

Vanishing of the Cosmological Constant, Stability of the Dilaton and Inflation

A. de la Macorra¹

*Instituto de Fisica UNAM, Apdo. Postal 20-364,
01000 Mexico D.F., Mexico.*

Abstract

We study the possibility of canceling the cosmological constant in supergravity string models. We show that with a suitable choice of superpotential the vacuum energy may vanish with the dilaton field at its minimum and supersymmetry broken with a large hierarchy. We derive the condition for which the introduction of a chiral potential, e.g. the inflaton potential, does not destabilize the dilaton field even in the region where the scalar potential takes positive values. This allows for an inflationary potential with the dilaton frozen at its minimum.

¹E-mail address: macorra@teorica0.ifisicacu.unam.mx

The vanishing of the cosmological constant remains an open issue [1]. In the absence of an understanding of why the cosmological constant cancels one can study, nevertheless, the possibility of having a potential that vanishes at the global minimum. In non-supersymmetric models one has, in principle, a large freedom to choose arbitrary terms and fine tune them to get zero cosmological constant. However, in a supersymmetric theory the potential terms are more constraint. In fact, for global supersymmetry (SUSY) this is not possible if one requires SUSY to be spontaneously broken. On the other hand, in supergravity models one can have the potential $V = 0$ and SUSY spontaneously broken. The breaking of SUSY is certainly a necessary condition but usually spontaneously broken symmetries have a potential with $V|_{min} = -\Lambda^4$, where Λ is the symmetry breaking scale and for a realistic hierarchy solution $V \simeq -[10^{-6} m_{Planck}]^4$, many orders of magnitude larger than the observational upper limit $|V| < 10^{-120} m_{Planck}^4$. It is therefore interesting to determine whether one can have vanishing cosmological constant with SUSY broken at a phenomenological acceptable scale in supergravity models.

The cosmological constant also plays an important role in the evolution of the universe. Since string models [2] are valid below the Planck scale they should describe the evolution of the universe. The standard big bang theory has some shortcomings like the isotropy and flatness problems [3]. An inflationary epoch [4], where the universe expanded in an accelerated way, may solve these problems. This inflationary period is also welcome to erase the abundance of the topological defects produced when a symmetry is spontaneously broken. In order for the potential to inflate a necessary condition is to have a positive energy. For arbitrary values of the different fields one expects V to be positive and to evolve to its minimum. In this evolution one would hope for an inflationary period. However, it is difficult to obtain an inflationary potential in string models due to the dynamics of the dilaton field S [11],[14]. When the scalar potential V evolves to the minimum of the dilaton, the universe, keeping all other fields fixed, does not go through an inflationary period. At the minimum, SUSY is broken and for vanishing v.e.v. of the chiral fields, the vacuum energy is negative and of the order of Λ^4 . So, it seems that either we have a positive potential with the wrong dynamics due to dilaton field or when the dilaton is at its minimum we have a negative potential. We will show that it is indeed possible to have a positive inflationary potential and S frozen.

The effective $D = 4$ superstring model is given by an $N = 1$ supergravity theory [5] with at least two gauge singlet fields S and T as well as an unspecified number of gauge chiral matter superfields ϕ . We will consider only an overall moduli T . The v.e.v. of the dilaton field S gives the gauge coupling constant $g^{-2} = \text{Re}S$ at the string scale while the real part of the moduli fields $\text{Re}T = R_i^2$ the radius of the compactified dimension. The tree level scalar potential is given by [5]

$$V_0 = \frac{1}{4}e^K(G_a(K^{-1})^a_b G^b - 3|W|^2) \quad (1)$$

where the Kahler potential is $G_a \equiv K_a W + W_a$ and

$$\begin{aligned} K &= -\log(S_r) - 3\log(T_r) + T_r^n |\phi|^2 \\ W &= W_0(S, T) + W_{ch}(T, \phi). \end{aligned} \quad (2)$$

The indices a, b run over all chiral fields, i.e. the dilaton S , the moduli T and chiral fields ϕ_i with modular weight n and $T_r = T + \bar{T}$, $S_r = S + \bar{S}$. In orbifold compactification n can vary from 5 to -9 [6]. All the fields are normalized with respect to the reduced Planck mass $m_p = M_p/\sqrt{8\pi}$. As usual, the indices a, b of the functions G, K and W denote derivatives with respect to chiral fields.

The superpotential term $W_0(S, T)$ arises due to non-perturbative effects, like gaugino condensation [7], and is responsible for breaking supersymmetry (SUSY) while W_{ch} is the chiral matter superpotential. Another important property of string models is duality invariance. Under this symmetry the moduli field T transform as element of $SL(2, Z)$ group, the dilaton S is invariant² while the chiral fields ϕ transform as modular function with weight n . The superpotential must be a modular function of weight -3. Using the duality symmetry one can then fix the T dependent part of the superpotential W in terms of the Dedekind-eta function η ,

$$W = \eta(T)^{-6} \Omega(S, \phi), \quad \eta \equiv e^{-\pi T/12} \Pi_n(1 - e^{-n\pi T}) \quad (3)$$

where we have assumed, for simplicity, that the T dependent part in W can be factorized and

$$\Omega = \Omega_0(S) + \Omega_{ch}(\phi). \quad (4)$$

Ω_0 corresponds to the non-perturbative superpotential due to gaugino condensate while Ω_{ch} to the chiral matter superpotential. Using eqs.(1) and (4) the scalar potential becomes

$$V_0 = \frac{1}{4} e^K |\eta|^{-12} \left[|h|^2 + |k_i|^2 + |\Omega|^2 \left(\frac{3T_r^2}{4\pi^2} |\hat{G}_2(T)|^2 - 3 \right) \right] \quad (5)$$

where \hat{G}_2 is the Eisenstein function of modular weight 2 and we have defined the auxiliary fields of S and ϕ in terms of $h = \eta^6 S_r G_S = S_r \Omega_S - \Omega$ and $k_i = \eta^6 G_i = K_i \Omega + \Omega_i$ respectively while the auxiliary field of the moduli T is $G_T = \frac{3}{2\pi} \eta^{-6} \Omega \hat{G}_2(T)$. Note that we have taken the chiral fields in eq.(3) to be canonically normalized, i.e. $K_\phi^\phi = 1$.

In the context of supergravity (tree level and one-loop level potential only), the cancelation of the cosmological constant must come through a non-vanishing auxiliary field $G_i \neq 0$. In string models there is a large number of chiral matter fields ϕ with different interacting superpotential terms. Since we do not wish to specify to any given model we will take Ω_{ch} to be arbitrary and we will study if there exist a particular superpotential for which SUSY is broken with the dilaton at its minimum and zero cosmological constant.

The condition of zero cosmological constant, considering only the tree level potential, is

$$G_a (K^{-1})^a_b G^b = 3|W|^2 \quad (6)$$

and it is hard to satisfy dynamically in the context of string models. In fact, the stabilization of the dilaton field is not simple to achieve and there are two different

²If loop corrections are included there may be a mixing term between S and T in the Kahler potential K and S would no longer be invariant. However, this loop correction is expected to be small with respect to the tree level value and we will for simplicity not consider it.

approaches that give a large hierarchy³. The first case is to consider two different gaugino condensate [8] with slightly different one-loop β -function coefficients and chiral matter fields with non-vanishing v.e.v. In this case a stable solution is found for $G_S = 0$. Another possibility is to include one-loop corrections of the 4-Gaugino interaction and to minimize $V_0 + V_1$ [9], where V_1 is the one-loop potential. In this case a stable minimum is found for a single gaugino condensate. In either case the effective S -dependent superpotential generated by a gaugino condensate is $\Omega_0(S) = d e^{-3S/2b_0}$ with b_0 the one-loop beta function coefficient and d independent of S .

I) Two Gaugino Condensates

Let us first consider the case of two gaugino condensates. To find the vacuum state with zero cosmological constant one needs to solve the eqs.

$$V| = V_S| = V_T| = V_i| = 0. \quad (7)$$

The extremum eq. for the moduli field T is

$$V_T = (K_T - 6\frac{\eta_T}{\eta})V + \frac{\partial}{\partial T} \frac{1}{4} e^K |\Omega|^2 \left[\frac{3T_r^2}{4\pi^2} |\hat{G}_2(T)|^2 \right] = 0 \quad (8)$$

and using $V| = 0$, eq.(8) is satisfied for T at the dual invariant points ($T = 1, e^{-\pi/6}$) where $\hat{G}_2 = 0$. This implies that the auxiliary field of the moduli is zero, $G_T = 0$, and it does not break SUSY contrary to the case where the condition $V| = 0$ was not imposed. For T at the dual points the vanishing of V implies then that either h or k or both must be different than zero.

The extremum eq. for the dilaton field is

$$V_S = \frac{1}{4} e^K |\eta|^{-12} \left[h_S \bar{h} + h \bar{h}_S + k_S \bar{k} + k \bar{k}_S - \frac{1}{S_r} (|h|^2 + |k|^2 - 3h\bar{\Omega}) \right] = 0 \quad (9)$$

which implies that

$$h_S \bar{h} + h \bar{h}_S - \frac{h}{S_r} (\bar{h} - 3\bar{\Omega}) = -\frac{\bar{k}}{S_r} (S_r k_S - k) \quad (10)$$

where we have already taken $\bar{k}_S = 0$. From eq.(10) we see that $h = 0$ is valid only if $\bar{k}(S_r k_S - k) = 0$. In [8] it was shown that for $k = 0$, $h = 0$ corresponds to a minimum while $h \neq 0$ to a maximum. However, imposing the condition $V| = 0$ we can no longer have $k = h = 0$. Assuming then that $k \neq 0$, $V_S = h = 0$ would require

$$S_r k_S - k = K_i (S_r \Omega_S - \Omega) + S_r \Omega_{S_i} - \Omega_i = 0. \quad (11)$$

The first term of eq.(11) vanishes by hypothesis ($h = 0$) and taking $\Omega_{S_i} = 0$ implies that $\Omega_i = 0$. Since $|k|^2 = |K_i \Omega + \Omega_i|^2 = |K_i \Omega|^2$ must cancel the term $3|\Omega|^2$ to have zero cosmological constant, a solution to $V| = V_S| = h = 0$ is valid only if

$$\Omega_i = 0 \quad \text{and} \quad K_i = \bar{\phi}_i = \sqrt{3}. \quad (12)$$

³We will not consider here the S -dual symmetry [13]

However, it is clear that this value cannot be dynamically obtained since the direction $\phi \rightarrow 0$ gives $k = \bar{\phi}\Omega \rightarrow 0$ and a negative V . We come to the conclusion that we cannot have $V| = V_S| = 0$ and $h = 0$. A different solution to S must be found in the context of two gaugino condensates. But, as we will show below, the new minimum for S requiring that $V| = 0$ is not far away from the one obtained for $h = 0$ with $V| < 0$.

Can we actually have $k \neq 0$? If all superpotential terms are at least quadratic in ϕ_i then $k = 0$ for $\phi_i = 0$. The only possibility to have $k \neq 0$ is with a linear superpotential $\Omega_{ch} = c\phi$, where c is an arbitrary constant. Let us take then the superpotential

$$\Omega = \Omega_0(S) + c\phi. \quad (13)$$

The auxiliary field $k = \bar{\phi}(\Omega_0 + c\phi) + c$ may vanish only if

$$|\Omega_0|^2 - 4|c|^2 \geq 0. \quad (14)$$

Therefore, for c large enough k will be different than zero. The scalar potential which is generated by the superpotential (13) is

$$V_0 = \frac{1}{4}e^K |\eta|^{-12} \left[|S_r \Omega_S - (\Omega_0 + c\phi)|^2 + |\bar{\phi}(\Omega_0 + c\phi) + c|^2 - 3|\Omega_0 + c\phi|^2 \right]. \quad (15)$$

In the case of two gaugino condensates Ω_0 is

$$\Omega_0 = d_1 e^{-3S/2b_1} + d_2 e^{-3S/2b_2} \quad (16)$$

where $b_i = \frac{3N_i - M_i}{16\pi^2}$ is the one-loop beta function coefficient of the i -th gauge groups and $d_i = \left(\frac{M_i}{3} - N_i\right) (32\pi^2 e)^{\frac{3(M_i - N_i)}{3N_i - M_i}} \left(\frac{M_i}{3}\right)^{\frac{M_i}{3N_i - M_i}}$ [8]. The eqs. $V| = V_S| = V_\phi = 0$ are

$$\begin{aligned} |h|^2 + |k|^2 &= 3|\Omega|^2 \\ h_S \bar{h} + h \bar{h}_S &= \Omega_S (3\bar{\Omega} - \bar{\phi}\bar{k}) \\ k_\phi \bar{k} + k \bar{k}_\phi &= \Omega_\phi (3\bar{\Omega} + \bar{h}) \end{aligned} \quad (17)$$

with $h_S = S_r \Omega_{SS}$, $\bar{h}_S = S_r \bar{\Omega}$, $k_\phi = \bar{\phi}\Omega_\phi$ and $\bar{k}_\phi = \bar{\Omega}$. One cannot solve the eqs. (17) analytically and the numerical solution yields

$$\langle S \rangle = 2.15, \quad \langle \phi \rangle = -0.73, \quad c = 2 \times 10^{-15} \quad (18)$$

where we have taken as an example $N_1 = 6, M_1 = 0, N_2 = 7, M_2 = 6$. This solution corresponds to a minimum.

More general, a solution is a minimum if the Hessian $H = \det V_{ab}$, where a, b denote derivatives with respect to real scalar fields and \det the determinant, is positive definite. Let us define $p = \text{Re } S$, $q = \text{Im } S$, $\phi = \rho e^{i\theta}$ with $\rho > 0$. If we only consider the contribution of the two gaugino condensates, i.e. there is no contribution of the chiral matter field ϕ , one has $H = V_{pp}V_{qq} - V_{pq}^2$. Since $q = \text{Im } S$ is an angular variable there will always be a minimum in the q direction (i.e. $V_{qq} > 0$). In fact, the q dependent part of the potential is $V(q) = A \cos[q(3/2b_1 + 3/2b_2)]$ with A independent of q (for eq.(15) we find $A = 2d_1 d_2 [p(3/b_1 + 3/b_2) - 1] e^{-p(3/2b_1 + 3/2b_2)}$). For $V_q = 0$ we necessarily have $V_{qp} = 0$ and therefore the potential has a minimum in the $\text{Re } S, \text{Im } S$

plane if and only if $V_{pp} > 0$. It is easy to see that $h = 0$ corresponds to $V_{pp} > 0$ while $h \neq 0$ to $V_{pp} < 0$. In this case $\langle S \rangle \simeq 2.16$ and we see that the v.e.v. of the dilaton is not very sensitive to the cancelation of the cosmological constant. Using eq.(18) the difference on $\langle S \rangle$ with or without the condition $V| = 0$ is $\Delta S = O(10^{-2})$.

Let us now consider the Hessian including the contribution of ϕ ,

$$H = \begin{vmatrix} V_{pp} & 0 & V_{\rho p} & 0 \\ 0 & V_{qq} & 0 & V_{\theta q} \\ V_{p\rho} & 0 & V_{\rho\rho} & 0 \\ 0 & V_{q\theta} & 0 & V_{\theta\theta} \end{vmatrix} \quad (19)$$

where again $V_{qp} = V_{q\rho} = V_{\theta p} = V_{\theta\rho} = 0$ for $V_q = V_\theta = 0$. The Hessian is positive definite if and only if $V_{pp}, V_{qq}, V_{\theta\theta}, V_{\rho\rho} > 0$ and $H_1 \equiv V_{pp}V_{qq} - V_{pq}^2 = V_{pp}V_{qq} > 0$, $H_2 \equiv V_{pp}V_{\rho\rho} - V_{p\rho}^2 > 0$, $H_3 \equiv V_{qq}V_{\theta\theta} - V_{q\theta}^2 > 0$. Since θ and q are angular variables there will always be a minimum in those directions and in the θ, q plane, i.e. $H_3 > 0$. An extremum solution will be a minimum if and only if $V_{pp}, V_{\rho\rho}, H_2 > 0$. For the example given above we have checked numerically that $V_{pp} > 0$, $V_{\rho\rho} > 0$ and $H_2 = V_{pp}V_{\rho\rho} - V_{p\rho}^2 > 0$ corresponding to a minimum. In [11] it was argued that a potential with $V| = V_S| = 0$ always has a global minimum with $V < 0$. However, their conclusion assumed that $V| = V_S| = 0$ with $W| = W_S| = 0$ which is not our case since $W, W_S \neq 0$ at the minimum. In **fig.1** we show V_0 as a function of S .

We have thus shown that it is possible to cancel the cosmological constant using only the tree level sugra scalar potential. The dilaton field has a stable solution and its v.e.v. does not differ much from the case where $\Omega_{ch} = 0$ and $V_{min} < 0$. SUSY is broken with a large hierarchy but mainly due to the auxiliary field $k = \eta^6(T) G_\phi \sim \Omega_\phi$ since $G_T = 0$ and $h = \eta^6(T) S_r G_s \sim \frac{\Omega_\phi \Omega_S}{\Omega_{SS}}$. Note that

$$\left| \frac{k}{h} \right| \simeq \left| \frac{\Omega_\phi \Omega_{SS}}{\Omega \Omega_S} \right| = O(10^3) \gg 1 \quad (20)$$

since $\Omega_{SS} \gg \Omega_S \sim \Omega$ and $\Omega_\phi \sim \Omega$.

A drop back is that many phenomenological quantities depend strongly on how SUSY is broken and in this case it is broken via the term which we now least and was introduced with the only motivation of rendering $V| = 0$.

II) One gaugino condensate

We will now study the scalar potential including loop corrections. The dilaton field is stabilized with a single gaugino condensate [9]. The Coleman-Weinberg one-loop potential is [10]

$$V_1 = \frac{\Lambda^4}{64\pi^2} \text{Str} \left[x + x^2 \ln\left(\frac{x}{1+x}\right) + \ln(1+x) \right] \quad (21)$$

where $x = \frac{M^2}{\Lambda_c^2}$ and M^2 are the (mass)² matrices for the different particles. The leading contribution to eq.(21) comes from the gaugino mass $m_g = \frac{b_0}{6\Lambda_c} e^{K/2} H$ with

$$H \equiv \eta^{-6} \left[F^S (K^{-1})_S^S (K_S \Omega + \Omega_S) + K^i (K^{-1})_i^i (K_i \Omega + \Omega_i) - 3\Omega \right] \quad (22)$$

and $F_S = -(1 + \frac{3S_r}{2b_0})$. Since $x_g = \frac{m_g^2}{\Lambda_c^2} = O(10^{-2})$ we can write eq.(21) as $V_1 = -\frac{n_g}{16\pi^2} \Lambda_c^2 m_g^2$ where Λ_c is the condensation scale and n_g is the dimension of the hidden

gauge group responsible for breaking SUSY. The scalar potential $V = V_0 + V_1$ can then be easily written as $V = V_0 - \frac{n_g}{16\pi^2} \Lambda_c^2 m_g^2$

$$\begin{aligned} V &= \frac{1}{4} e^K \eta^{-12} \left[|h|^2 + |k|^2 - 3|\Omega|^2 - \delta |H|^2 \right] \\ V &= \frac{1}{4} e^K \eta^{-12} \left[|h|^2 (1 - \delta F_S^2) + |k|^2 (1 - \delta |\phi|^2) - 3|\Omega|^2 (1 + 3\delta) \right. \\ &\quad \left. - \delta (F_S h [\phi \bar{k} - 3\Omega] - 3\phi \Omega k + h.c.) \right] \end{aligned} \quad (23)$$

where we have defined $\delta \equiv \frac{n_g b_0^2}{144\pi^2} \ll 1$ and we have specialized to $K_i = \bar{\phi}$. We recover the tree level potential by setting $\delta = 0$. In the absence of chiral matter potential, i.e. $k = 0$, the dominant term in eq.(23) is

$$|h|^2 (1 - \delta F_S^2) = \left| de^{-3S/2b_0} \right|^2 \left(1 + \frac{3S_r}{2b_0} \right) \left(1 - \delta \left[1 + \frac{3S_r}{2b_0} \right]^2 \right) \quad (24)$$

where we have taken a single gaugino condensate $\Omega_0 = de^{-3S/2b_0}$ for which $h = -\Omega_0 \left(1 + \frac{3S_r}{2b_0} \right)$. It is clear from eqs.(23) and (24) that there is stable solution to S with a single gaugino condensate unlike the previous case (tree level only). Using this approximation, the minimum is at

$$S_r \simeq \frac{2b_0}{3\delta^{1/2}} = \frac{8\pi}{\sqrt{n_g}} \quad (25)$$

and

$$(1 - \delta F_S^2) \simeq -2\delta^{-1/2} < 0. \quad (26)$$

For $k \neq 0$ a vanishing cosmological constant requires

$$|k|^2 = -|h|^2 (1 - \delta F_S^2) + 3|\Omega|^2 \quad (27)$$

up to leading order. In this case SUSY is broken mainly due to the auxiliary field of the dilaton since

$$\frac{|k|^2}{|h|^2} \simeq -(1 - \delta F_S^2) = 2\delta^{1/2} \sim 10^{-2}. \quad (28)$$

Solving the exact eqs $V = V_S = V_\phi = 0$ numerically, using the ex. 1 in [9], i.e. $SU(6)$ gauge group with $b_0 = 11/16\pi^2$, the v.e.v. of S and ϕ and the value of c are

$$S = 2.11, \quad \phi = 0.74, \quad c = 2.9 \times 10^{-19} \quad (29)$$

corresponding to a stable solution, i.e. $V_{pp} > 0$, $V_{\rho\rho} > 0$ and $H_2 = V_{pp}V_{\rho\rho} - V_{p\rho}^2 > 0$. The difference between the v.e.v. of S when $c = \phi = 0$ and $V < 0$ and when the condition $V| = 0$ is imposed is small $\Delta S = O(10^{-3})$.

In **fig.2** we show V as a function of S and in **fig.3** we have V vs (ϕ, S) . We see that V has two minimums in the S direction. In the absence of the chiral matter field contribution, i.e. $c = \phi = 0$, V has only one minimum and a similar behavior as the one shown in **fig.1**. However, for $c\phi \neq 0$ there is a local minimum with $V > 0$. Let us call S_0 the global minimum, S_2 the local minimum and S_1 the maximum in between both minima. In the region $S < S_0$ the potential V_0 and $(-V_1)$ decrease

exponentially fast and we have $S_r \Omega_S = \frac{3S_r}{2b_0} \Omega_0 \gg c\phi$ and $x_g \ll 1$. The dominant term in V is $|h|^2(1 - \delta F_S^2) + |k|^2 \sim |S_r \Omega_S|^2(1 - \delta F_S^2) + |c|^2$ and the minimum is at S_0 . For $S > S_1$ the tree level potential decreases as $1/S$ since in this region $c\phi > \Omega_S$ and the dominant term in V_0 is then $e^K |c|^2 \sim |c|^2/S$. On the other hand, the one-loop potential V_1 decreases slightly (it becomes more negative) in the region $S_1 < S < S_2$ and therefore we have a local minima at S_2 . This behavior of $V = V_0 + V_1$ can be seen from eq.(23) since the first two terms dominate and if S becomes larger then $|h|^2(1 - \delta F_S^2) + |k|^2 \sim -|c\phi|^2 \delta \left(\frac{3S_r}{2b_0}\right)^2 + |c|^2$ decreases. For $S > S_2$ the "log" terms in V_1 (cf. eq.(21)) start to dominate, $x_g \gg 1$ and $V_1 \sim -\Lambda_c^4 \ln x_g$ decreasing exponentially fast.

Again, we have seen that it is possible to cancel the cosmological constant while having S at its minimum with a large hierarchy.

We will now consider the inflationary potential. The dynamics of the dilaton field does not allow for an inflationary potential as long as S is still a dynamically field. A possible solution is to have S froze at its minimum and then inflation could be driven by other fields. To be more general, we will investigate the potential due to chiral matter fields other than ϕ which we used to cancel the cosmological constant⁴. Clearly, a potential to inflate must be positive. Is it then possible to have S stable and $V > 0$? Can the inflaton potential dominate V ?

Let us take the potential V to be independent of the inflaton field, which we will denote by Φ . We assume V to have the properties $V|_{S_0} = V_S|_{S_0} = 0$ at the global minima which we have seen that it is indeed feasible. The inflaton potential $V_{inf}(S, \Phi)$ will in general depend on S . It is reasonable to assume that the S dependent part of V_{inf} can be factorized so that

$$V_{inf} = A(S)P(\Phi). \quad (30)$$

For the usual inflaton potential in string scenarios one can have $A = 1/S \sim e^K$ or $A = e^{-3S/2b_0}/S \sim e^K W_0$ dependence on the dilaton field. The complete scalar potential is then

$$V_f = V + V_{inf}. \quad (31)$$

Due to the inclusion of V_{inf} the value of S at the minimum will now vary and we will call it S_1

$$V_{f,S}|_{S_1} = V_S|_{S_1} + V_{inf,S}|_{S_1} = 0 \quad (32)$$

which implies that

$$P(\Phi)A_S|_{S_1} = -V_S|_{S_1}. \quad (33)$$

Expanding V and V_S around S_0 and keeping the leading term only, one has

$$\begin{aligned} V &= \frac{1}{2}V_{SS}|_{S_0}(S - S_0)^2 + \dots \\ V_S &= V_{SS}|_{S_0}(S - S_0) + \dots \end{aligned} \quad (34)$$

since $V|_{S_0} = V_S|_{S_0} = 0$ by hypothesis and ... correspond to higher derivative terms.

⁴It would, nevertheless, be interesting to study the phenomenology of ϕ .

Let us take $A_S = \alpha(S)A(S)$. For $A = 1/S$ one has $\alpha = -1/S$ while for $A = e^{-3S/2b_0}/S$ one has $\alpha = -3/2b_0$ (keeping the leading contribution only since $3/2b_0 \gg 1$). Using eqs.(32)-(34) we can write the inflaton potential $V_{inf} = A(S)P(\Phi)$ as

$$\begin{aligned} V_{inf}(S_1) &= \frac{1}{\alpha} A_S(S) P(\Phi) \\ V_{inf}(S_1) &\simeq -\frac{1}{\alpha} V_{SS|S_0} (S_1 - S_0) \end{aligned} \quad (35)$$

and the complete scalar potential is

$$V_f(S_1) = V_{SS|S_0} (S_1 - S_0) \left(\frac{1}{2} (S_1 - S_0) - \frac{1}{\alpha} \right). \quad (36)$$

The first term in eq.(36) corresponds to the contribution from V while the second comes from V_{inf} . Which terms dominates depends on $S_1 - S_0$ and α . For $\alpha(S_1) = -1/S_1$ the potential is $V_{inf} = V_{SS|S_0} (S_1 - S_0) \left(\frac{1}{2} (S_1 - S_0) + S_1 \right)$ and clearly the second term, i.e. V_{inf} dominates. For $\alpha = -3/2b_0$ the potential is $V_{inf} = V_{SS|S_0} (S_1 - S_0) \left(\frac{1}{2} (S_1 - S_0) + 2b_0/3 \right)$ and V_{inf} will dominate only if $(S_1 - S_0) < 4b_0/3$. Note that eqs.(34) and (35) are only valid if S_1 does not differ much from S_0 and it must be calculated for each specific example. Furthermore, the condition on $S_1 - S_0$ to be small, i.e. that V_{inf} does not destabilize S , imposes a constraint on the magnitude of V_{inf} [14]. Using eqs.(35) and (36) the inflaton potential must satisfy

$$\left| \frac{V_{inf}}{V_{SS|S_0}} \right| < \frac{2}{\alpha^2}. \quad (37)$$

This condition has been used to set an upper limit to the density fluctuations [14] and for specific examples it can be consistent with the COBE observations [12].

To conclude, we have shown that in the context of sugra string models the cosmological constant can be arranged to vanish if the superpotential has a linear term. The solution corresponds to a global minimum. We have also shown that the introduction of a chiral potential V_{inf} does not necessarily destabilize the dilaton field and the vacuum energy can nevertheless be dominated by V_{inf} allowing for the existence of an inflationary potential.

Acknowledgment

It is a pleasure to thank G. G. Ross for suggesting the problem and for many enlightening discussions and comments.

References

- [1] For a review of the Standard Cosmological Model see S. Weinberg "Gravitation and Cosmology Principles and Applications of the General Theory of Relativity", J. Wiley, New York (1972).
- [2] For a review of string theories, see M. Green, J. Schwarz and E. Witten, Superstring Theory, Cambridge University Press, 1987.

- [3] For a review of the open questions of the standard hot big-bang theory see E. W. Kolb and M. S. Turner “*The Early Universe*”, Frontiers in Physics 69 (1990), chapter 8 and references therein.
- [4] A. D. Linde, Phys. Lett. B129, 177 (1983) and ref. therein.
- [5] E. Cremmer, S. Ferrara, L. Girardello and A. Van Proeyen, Nucl.Phys. B212 (1983) 413.
- [6] L.E. Ibanez and D. Lust, Nucl. Phys. B382 (1992) 305.
- [7] For a review see D. Amati, K. Konishi, Y. Meurice, G. Rossi and G. Veneziano, Phys. Rep 162 (1988) 169; J.P.Derendinger, L.E.Ibanez and H.P.Nilles, Phys. Lett. B155 (1985) 65; M.Dine, R.Rohm, N.Seiberg and E.Witten, Phys. Lett. B156 (1985) 55; A. Font, L. Ibanez, D. Lust and F. Quevedo, Phys. Lett. B245 (1990) 401.
- [8] B. de Carlos, J. A. Casas and C. Munoz, Phys. Lett. B263 (1991) 248.
- [9] A. De La Macorra and G. G. Ross, Nucl. Phys. B404 (1993) 321; Phys. Lett. B325 (1994) 85.
- [10] S. Coleman and E. Weinberg, Phys. Rev. D7 (1973) 1883.
- [11] R. Brustein and P. J. Steinhardt, Phys. Lett. B302 (1993) 196;
- [12] COBE observations of the CMBR, Physics Today, June 1992.
- [13] A. Font, L. E. Ibanez, D. Lüst and F. Quevedo, Phys. Lett. B249 (1990) 35.
- [14] A. de la Macorra and S. Lola, *Inflation from Superstrings*, hep-ph/9411443, UNAM preprint IFUNAM-FT-94-63, Heidelber preprint HD-THEP-94-45.

Fig.1 We show the potential of two gaugino condensates V_0 as a function of S .

Fig.2 We show the potential of a single gaugino condensate $V = V_0 + V_1$ as a function of S for $c\phi \neq 0$.

Fig.3 We show $V = V_0 + V_1$ as a function of S and ϕ .

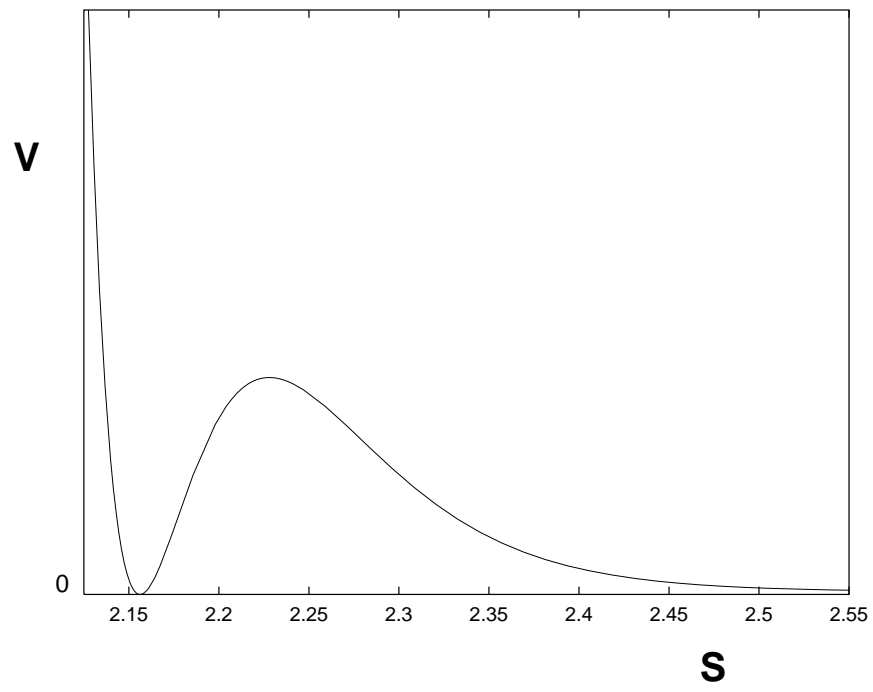


Fig.1

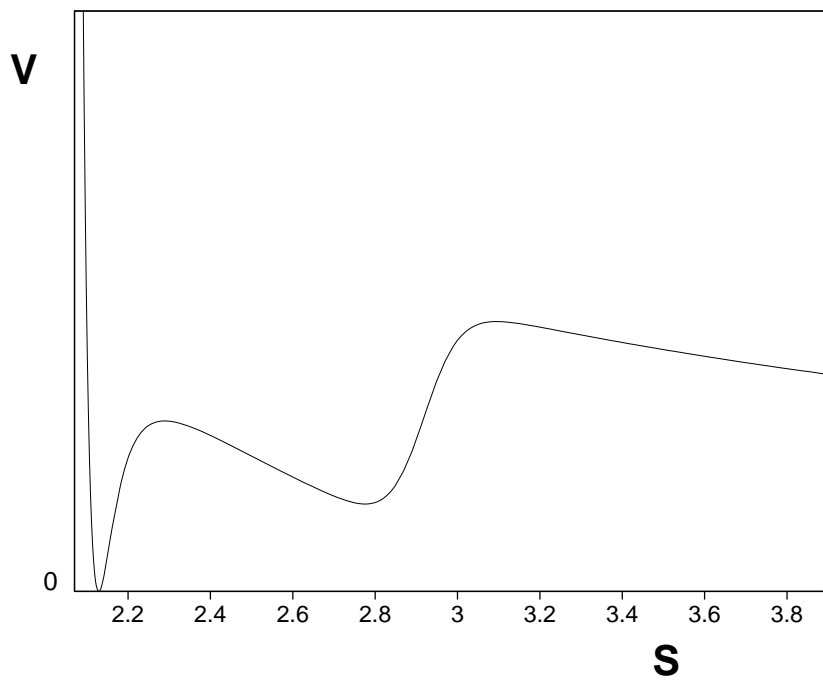


Fig.2

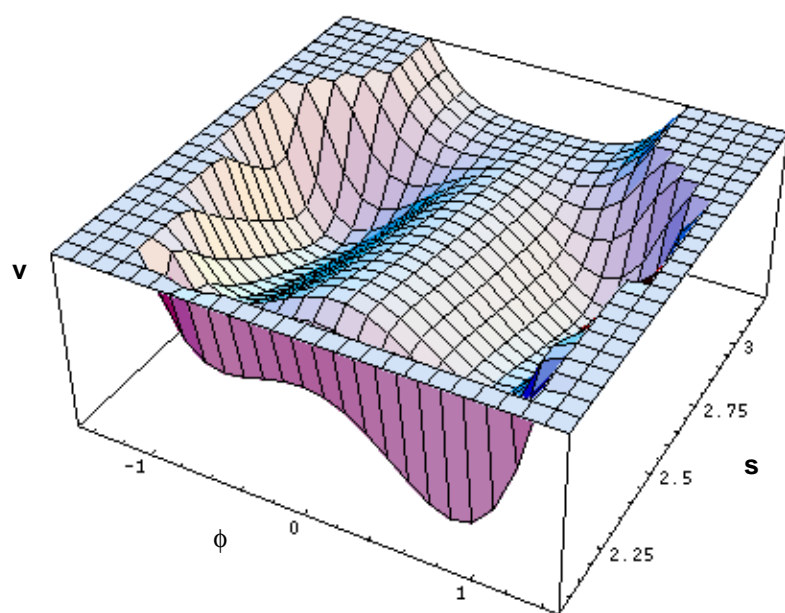


Fig.3